

HYPERONS IN THE BOUND STATE APPROACH WITH VECTOR MESONS

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December 30, 1993

ABSTRACT

We investigate a model for hyperons based on the bound state approach in which vector mesons are explicitly incorporated. We show that for empirical values of the mesonic parameters the strange hyperon spectrum is well reproduced. We also discuss the extension of the model to heavier flavors. We show that the explicit presence of the heavy vectors leads to good predictions for the heavy baryon masses.

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The bound state approach [1] has become a very useful tool for the study of hyperon properties [2]. In this approach strange hyperons are described as soliton-kaon bound systems. To find the soliton-kaon interactions one usually starts with an effective chiral lagrangian supplemented with proper chiral symmetry breaking terms. Up to now, most of the calculations have been done using effective lagrangians where only pseudoscalar meson fields were included. There are several reasons to believe, however, that these are not the only degrees of freedom to be taken into account. Within the chiral soliton model vector mesons are known to improve the description of the nucleon properties [3]. Moreover, it has been recently argued [4, 5] that they are of fundamental importance in the extension of the bound state description to heavier flavors. It is now well-known that in the heavy quark limit of QCD the heavy vector meson (i.e. B^*) becomes degenerate with the heavy pseudoscalar (i.e. B) and therefore should be explicitly included in the effective action [6]. In the past, some attempts have been done to explicitly incorporate vector mesons in the bound state model. In Ref.[7] a model with an explicit ω vector meson was studied. In Ref.[8] although the full nonet of vector mesons was included, the strange vector meson (K^*) was integrated out before solving the equations of motion. In the present paper we will deal with the full vector meson nonet throughout. In addition, a more general chiral symmetry breaking action that includes derivative-type terms will be used. Finally, although most of this work will be devoted to the description of low-lying positive parity strange hyperons some results for charmed and bottom baryons will also be given.

To incorporate vector mesons in the chiral lagrangian we use the hidden gauge approach [9]. This scheme is based on the identification of the “strong” vector mesons as gauge bosons of the hidden gauge symmetry of the non-linear sigma model. The corresponding effective action is given by

$$\Gamma = \Gamma_0 + \Gamma_{sb} + \Gamma_{an}, \quad (1)$$

where Γ_0 is the non-anomalous chirally symmetric contribution to the action

$$\begin{aligned} \Gamma_0 = & \int d^4x \left\{ \frac{1}{2g^2} \text{Tr } F_{\mu\nu}F^{\mu\nu} \right. \\ & \left. - \frac{f_\pi^2}{4} \left[\text{Tr} \left(\mathcal{D}_\mu \xi_L \xi_L^\dagger - \mathcal{D}_\mu \xi_R \xi_R^\dagger \right)^2 + a \text{Tr} \left(\mathcal{D}_\mu \xi_L \xi_L^\dagger + \mathcal{D}_\mu \xi_R \xi_R^\dagger \right)^2 \right] \right\}. \end{aligned} \quad (2)$$

The covariant derivatives are

$$\mathcal{D}_\mu \xi_{L(R)} = (\partial_\mu - V_\mu) \xi_{L(R)} \quad (3)$$

and the field strength tensor is

$$F_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu - [V_\mu, V_\nu]. \quad (4)$$

The antihermitian vector meson field V_μ is defined as

$$V_\mu = \frac{1}{2} \begin{pmatrix} \omega_\mu + \rho_\mu & \sqrt{2}K_\mu^* \\ -\sqrt{2}K_\mu^{*\dagger} & \sqrt{2}\phi_\mu \end{pmatrix} \quad (5)$$

where ω_μ corresponds to the isoscalar vector meson, ρ_μ to the isovector vector meson, K_μ^* to the isodoublet of strange vector mesons and ϕ_μ to the isosinglet strange vector meson. Although the constant a in Eq.(2) is, strictly speaking, an arbitrary constant within the framework adopted here, the value $a = 2$ leads to a very successful meson phenomenology. In fact, using that value one obtains the KSRF relation as well as vector meson dominance[9]. In what follows, we will take the value $a = 2$ throughout.

In Eq.(1) we have included a chiral symmetry breaking term Γ_{sb} whose expression is given by [10]

$$\begin{aligned} \Gamma_{sb} = & \int d^4x \left\{ \frac{m_\pi^2 f_\pi^2 + 2m_K^2 f_K^2}{12} Tr[U - 1] + \frac{m_\pi^2 f_\pi^2 - m_K^2 f_K^2}{6} Tr[\sqrt{3}\lambda_8 U] \right. \\ & - Tr[(1 - \sqrt{3}\lambda_8)(\alpha (\mathcal{D}_\mu \xi_L)^\dagger \mathcal{D}^\mu \xi_R + \beta U^\dagger \partial_\mu U \partial^\mu U^\dagger + \gamma \xi_L^\dagger F_{\mu\nu} F^{\mu\nu} \xi_R)] \Big\} \\ & + h.c. \end{aligned} \quad (6)$$

Here, we have considered only those chiral symmetry breaking terms in Ref.[10] which are linear in the quark mass matrix and satisfy the OZI rule. In addition we have neglected the non-strange quark masses in the derivative-type symmetry breakers. These approximations do not affect in any significant way the properties of the mesons we are interested in. In Eq.(6), U is the chiral field ($= \xi_L^\dagger \xi_R$), m_π is the pion mass, m_K the kaon mass and f_π and f_K are the pion and kaon decay constants respectively. The coefficients α , β and γ are defined as follows

$$\alpha = \frac{1}{3g^2} (m_v^2 - Z_{K^*}^2 m_{K^*}^2) \quad (7)$$

$$\beta = \frac{1}{12} (f_\pi^2 - f_K^2 + 3\alpha) \quad (8)$$

$$\gamma = \frac{1}{6g^2} (1 - Z_{K^*}^2) \quad (9)$$

where m_v ($= \sqrt{a}g f_\pi$) is the non-strange vector meson mass, m_{K^*} is the K^* mass and Z_{K^*} is the renormalization constant of the K^* meson field.

Finally, the anomalous action Γ_{an} is given by

$$\Gamma_{an} = \Gamma_{WZ}^0 + \frac{iN_c}{24\pi^2} \int \left(Tr \left(a_L^3 a_R - a_R^3 a_L \right) - Tr \left(a_L a_R a_L a_R \right) \right) \quad (10)$$

where Γ_{WZ}^0 is the irreducible Wess-Zumino action containing only the pseudoscalar fields and $a_{L(R)} = \mathcal{D}_\mu \xi_{L(R)} \xi_{L(R)}^\dagger dx^\mu$. Eq.(10) is a particular combination of the general lowest order anomalous action given in Refs.[11, 12]. As done in Ref.[8] we will simplify its rather cumbersome form by expressing the anomalous ρ and K^* vector meson couplings in terms of the pseudoscalar fields. With this approximation our model turns out to be an $SU(3)$ generalization of the minimal model studied in Ref.[13]. No major difference in the predictions of the baryon observables has been found in models where the full anomalous action was included [12].

To study the baryon sector of the model we have to define an ansatz for the ξ_L and ξ_R fields. As in Ref.[8] we use

$$\begin{aligned} \xi_L^\dagger &= \sqrt{U_\pi} \sqrt{U_K} \\ \xi_R &= \sqrt{U_K} \sqrt{U_\pi} \end{aligned} \quad (11)$$

where

$$U_\pi = \begin{pmatrix} N^2 & 0 \\ 0 & 1 \end{pmatrix}, \quad N = \exp \left[i \frac{\vec{\tau} \cdot \vec{\pi}}{2f_\pi} \right], \quad (12)$$

$$U_K = \exp \left[\frac{i\sqrt{2}}{f_\pi} \begin{pmatrix} 0 & K \\ K^\dagger & 0 \end{pmatrix} \right], \quad K = \begin{pmatrix} K^+ \\ K^0 \end{pmatrix}.$$

Following the usual steps in the bound state approach we can now obtain the soliton-kaon action by inserting the ansatz Eq.(11) into the effective action and expanding up to second order in kaon fields. To make this expansion consistently, only terms up to second order in K^* should be retained. The resulting lagrangian density can be separated in two classes of terms. The first ones reduce to the lagrangian of interacting nonstrange fields (pions and ω and ρ vector mesons) which we will denote $\mathcal{L}_{SU(2)}$. The minimization of the corresponding action using the Skyrme-Wu-Yang ansatz provides the π , ω and ρ solitonic fields. Details of this procedure as well as the explicit expression of $\mathcal{L}_{SU(2)}$ can be found in Ref.[13]. The rest of the terms describe the coupled system of quantum K and K^* fields moving in the presence of the background fields. It is interesting to note that to quadratic order in K fields, the ϕ vector meson is completely decoupled from the KK^* system. Therefore its

contribution to \mathcal{L} will be ignored. In terms of the renormalized K and K^* fields¹ the resulting meson-soliton lagrangian reads

$$\begin{aligned}
\mathcal{L} = & \mathcal{L}_{SU(2)} + (D_\mu K)^\dagger D^\mu K + K^\dagger a_\mu a^\mu K - m_K^2 K^\dagger K \\
& + \frac{1}{2g^2 f_K^2} (3m_v^2 - 2Z_{K^*}^2 m_{K^*}^2) [K^\dagger (v_\mu + q_\mu) D^\mu K - (D^\mu K)^\dagger (v_\mu + q_\mu) K] \\
& - \frac{1}{g^2 f_K^2} K^\dagger \left[(2m_v^2 - Z_{K^*}^2 m_{K^*}^2) a_\mu a^\mu + (m_v^2 - Z_{K^*}^2 m_{K^*}^2) (q_\mu + v_\mu) (q^\mu + v^\mu) \right. \\
& \quad \left. - \frac{1 - Z_{K^*}^2}{2} q_{\mu\nu} q^{\mu\nu} + \frac{m_\pi^2 g^2 f_\pi^2}{4} (U_\pi + U_\pi^\dagger - 2) \right] K \\
& + i \frac{m_v^2}{g^2 f_K Z_{K^*}} [K^\dagger a_\mu K^{*\mu} + K_\mu^{*\dagger} a^\mu K] \\
& - \frac{1}{2g^2} \left[K_{\mu\nu}^{*\dagger} K^{*\mu\nu} - 2m_{K^*}^2 K_\mu^{*\dagger} K^{*\mu} + \frac{2}{Z_{K^*}^2} K_\mu^{*\dagger} q^{\mu\nu} K_\nu^* \right] \\
& + \frac{i N_c}{8f_K^2} \left[B_\mu \left((D^\mu K)^\dagger K - K^\dagger D^\mu K \right) \right. \\
& \quad \left. + 3\omega_\mu B^\mu K^\dagger K + \frac{4}{3\pi^2} \epsilon^{\mu\nu\beta\eta} \omega_\mu (D_\nu K)^\dagger a_\beta D_\eta K \right] \tag{13}
\end{aligned}$$

where

$$\begin{aligned}
D_\mu K &= \partial_\mu K + v_\mu K, \\
q_\mu &= \frac{1}{2}(\rho_\mu + \omega_\mu), \\
q_{\mu\nu} &= \partial_\mu q_\nu - \partial_\nu q_\mu - [q_\mu, q_\nu], \\
\begin{pmatrix} v_\mu \\ a_\mu \end{pmatrix} &= \frac{1}{2} (N^\dagger \partial_\mu N \pm N \partial_\mu N^\dagger), \\
B^\mu &= \frac{1}{24\pi^2} \epsilon^{\mu\nu\beta\eta} Tr \left[U_\pi^\dagger (\partial_\nu U_\pi) U_\pi^\dagger (\partial_\beta U_\pi) U_\pi^\dagger (\partial_\eta U_\pi) \right], \\
K_{\mu\nu}^* &= \partial_\mu K_\nu^* - \partial_\nu K_\mu^* - q_\mu K_\nu^* + q_\nu K_\mu^*. \tag{14}
\end{aligned}$$

To find the $\mathcal{O}(N_c^0)$ corrections to the hyperon properties we should now obtain the equations of motion of the K and K^* meson fields in the presence of the static soliton background. Since in this work we are only interested in the properties of the low-lying positive parity hyperons it is enough to find the ground state of the coupled

¹ To obtain the canonical form of the free action in the limit in which meson-soliton interactions are neglected the strange meson fields have to be renormalized. This is due to the presence of the derivative-type symmetry breakers in Eq.(6).

meson system. For that particular state the meson fields can be expressed as follows

$$\begin{aligned} K(\vec{r}, t) &= \frac{1}{\sqrt{4\pi}} k(r, t) \vec{\tau} \cdot \hat{r} \chi \\ K_0^*(\vec{r}, t) &= \frac{g}{\sqrt{4\pi}} t_0(r, t) \chi \\ \vec{K}^*(\vec{r}, t) &= \frac{g}{\sqrt{4\pi}} [t_2(r, t) \hat{r} - i t_3(r, t) \hat{r} \times \vec{\tau}] \chi \end{aligned} \quad (15)$$

where χ is the two-component isospinor. It is interesting to note that in the ground state the K^* vector meson field can be expressed in terms of only 3 time-dependent radial functions. In addition, as usual when dealing with vector fields, there is no dynamics associated with t_0 . As a consequence, the diagonalization of the corresponding meson hamiltonian leads, for the ground state, to a system of 3 coupled eigenvalue equations. Their lengthy expressions will be given elsewhere. The numerical solution of this eigenvalue problem provides the $\mathcal{O}(N_c^0)$ correction to the hyperon masses ϵ together with the meson eigenfunctions.

Finally, the splittings between hyperons with the same value of strangeness but different spin and/or isospin are obtained to $\mathcal{O}(N_c^{-1})$ by introducing the soliton $SU(2)$ collective rotations $A(t)$

$$\begin{aligned} U_\pi &\rightarrow AU_\pi A^\dagger \\ K &\rightarrow AK \\ K_\mu^* &\rightarrow AK_\mu^*. \end{aligned} \quad (16)$$

As it has been shown in Ref.[13] to this order some new components of the ρ and ω vector fields are excited. They contribute, together with the rest of the background fields, to the $SU(2)$ moment of inertia Θ . Details of the quantization of the $SU(2)$ sector of the present model can be found in Ref.[13]. The strange sector contributes to the rotational lagrangian with a term proportional to the scalar product of the induced meson spin times the soliton angular velocity [1]. The constant of proportionality is the so-called hyperfine splitting constant c . It can be calculated as the volume integral of some complicated function of the $SU(2)$ fields and the K and K^* eigenfunctions. Its explicit form will be also given elsewhere.

Summing up all the contributions to the meson-soliton hamiltonian and taking matrix elements between hyperon states one obtains the bound-state mass formula. It reads[1, 2]

$$M_{I,J,S} = M_{sol} + |S|\epsilon + \frac{1}{2\Theta} \left[cJ(J+1) + (1-c)I(I+1) + \frac{c(c-1)}{4}|S|(|S|+2) \right] \quad (17)$$

where I , J and S are the hyperon isospin, spin and strangeness, respectively.

In our numerical calculations we set all the parameters that appear in the effective lagrangian to their empirical values. Namely, for the parameters that determine the hadron properties in the $SU(2)$ sector we use $f_\pi = 93 \text{ MeV}$, $m_\pi = 138 \text{ MeV}$, $g = 5.85$ and $m_v = m_\rho = m_\omega = 770 \text{ MeV}$. Using these values the soliton mass turns out to be $M_{sol} = 1470 \text{ MeV}$ while the $SU(2)$ moment of inertia is $\Theta = 0.82 \text{ fm}$. Although the calculated moment of inertia compares fairly well with the best fit value $\Theta^{BF} = 1.01 \text{ fm}$ needed to reproduce the empirical $\Delta - N$ mass splitting, the soliton mass is much larger than $M_{sol}^{BF} = 866 \text{ MeV}$ needed to fit the nucleon mass. This large value of M_{sol} obtained for empirical f_π is an old problem in skyrmion physics. Quite recently it has been argued [14, 15] that this problem can be solved by taking into account the $\mathcal{O}(N_c^0)$ Casimir corrections due to pion loops. This might justify the traditional attitude we follow here of ignoring this problem by considering only hyperon masses taken with respect to the nucleon mass. For the mesonic parameters in the $SU(3)$ sector of the model we use $f_K = 114 \text{ MeV}$, $m_K = 495 \text{ MeV}$ and $m_{K^*} = 892 \text{ MeV}$. Within the present model, however, the value of Z_{K^*} cannot be consistently determined from the strange meson properties[10]. In fact, while the $K^* \rightarrow K\pi$ decay ratio favors the value $Z_{K^*} = 0.75$ from the empirical ϕ meson mass one obtains $Z_{K^*} \simeq 0.96$. For this reason we will consider Z_{K^*} as a free parameter within that range of values and study the behavior of the bound meson energy ϵ and the hyperfine splitting constant c as a function of it. Results are shown in Fig.1. We observe that ϵ depends quite strongly on Z_{K^*} . On the other hand, the constant c is less sensitive to Z_{K^*} . In the context of the Skyrme model a similar behavior has been found as a function of the kaon renormalization constant [16]. We also see that for $Z_{K^*} = 0.85$ the calculated values of $\epsilon = 217 \text{ MeV}$ and $c = 0.65$ are very close to the best fit values[2] $\epsilon^{BF} = 218 \text{ MeV}$ and $c^{BF} = 0.66$. In addition, this value of Z_{K^*} agrees very well with $Z_{K^*} = 0.84$ used in Ref.[10] to obtain a good overall fit of the meson properties ².

Before presenting our predictions for the hyperon masses some remarks are in order. First, we should stress the importance of the derivative-type symmetry breaking terms. Namely, if we set $\alpha = \beta = \gamma = 0$ in Eq.(6) we obtain $\epsilon = 110 \text{ MeV}$. As in the case of the Skyrme model strange mesons are overbound in the absence of such terms [16]. Second, when comparing this value with the results of Ref.[8] where these terms have been neglected and strange vector mesons have been completely integrated out we observe that the explicit presence of these mesons tends to increase the binding energy. In the present calculation this effect is compensated by the

²For $Z_{K^*} = 0.85$ we obtain $\Gamma(K^* \rightarrow K\pi) = 39 \text{ MeV}$ and $m_\phi = 1.12 \text{ GeV}$ to be compared with the empirical values 50 MeV and 1.02 GeV respectively.

derivative-type symmetry breakers.

In Table 1 we show our prediction for the strange hyperon mass spectrum. We observe that the model predicts the mass differences within few percent. The main source of disagreement is the somewhat small value of the calculated $SU(2)$ moment of inertia. It is interesting to note that if this moment of inertia is taken to its best fit value all the hyperon masses are predicted within roughly 10 MeV from their empirical values. It would be interesting to see whether effects like i.e. pion loop corrections could contribute to increase the value of the predicted $SU(2)$ moment of inertia.

Finally, we will briefly discuss the extension of the present model to the charm and bottom sectors. As already mentioned in those cases heavy vector mesons are expected to play a crucial role in the determination of heavy meson binding energies. In the charm sector, using $f_D = 158$ MeV and $Z_{D^*} = 0.47$, we predict $\Lambda_C - N = 1343$ MeV and $\Sigma_C - N = 1553$ MeV that compare very well with the empirical values 1346 MeV and 1514 MeV respectively. In the bottom sector, using $f_B = 149$ MeV and $Z_{B^*} = 0.32$ we predict $\Lambda_B - N = 4704$ MeV and $\Sigma_B^* - \Sigma_B \simeq 8$ MeV. The $\Lambda_B - N$ splitting agrees very nicely with the empirical value 4702 MeV while the hyperfine splitting $\Sigma_B^* - \Sigma_B$ is very close to zero as required by heavy quark symmetry. In all these calculations the heavy meson masses have been set to the empirical values given in Ref.[17]. From the results above it is clear that the overbinding problems found in previous calculations where only pseudoscalars were present [4] can be completely overcome in the present model without affecting the predictions for the hyperfine splittings. A very interesting point for further investigation is the connection between the present model and those based on explicit heavy quark symmetry [5, 18, 19].

In conclusion, we have studied a model for hyperons based on the bound state approach in which vector mesons are explicitly incorporated both in the light and massive sectors. We have shown that good predictions for the hyperon spectrum can be obtained for empirical values of the mesonic parameters. Derivative-type chiral symmetry breaking terms are crucial in obtaining these results. Finally, we have also shown that when extended to the charm and bottom sectors the present model eliminates the overbinding found in models where only pseudoscalars are present.

Discussions with E. Jenkins, A. Manohar, Y. Oh, D.O. Riska, M. Rho and K. Yamawaki are greatly acknowledged.

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Figure and Table Captions

Fig. 1: Meson eigenenergy ϵ and hyperfine splitting constant c as functions of the K^* renormalization constant Z_{K^*} .

Table 1: Strange hyperon masses (in MeV) taken with respect to the calculated nucleon mass. The experimental values are given in the first column. They correspond to the mass differences between isospin multiplet averages, with uncertainties given by half the sum of the difference of extremum values within each multiplet. The calculated values are given in the second column. The mesonic parameters used in this calculation are $f_\pi = 93 MeV$, $m_\pi = 138 MeV$, $f_K = 114 MeV$, $m_K = 495 MeV$, $g = 5.85$, $m_v = 770 MeV$, $m_{K^*} = 892 MeV$ and $Z_{K^*} = 0.85$.

Table 1

Particle	Exp.	This model
Λ	177 ± 1	165
Σ	254 ± 5	249
Σ^*	446 ± 3	484
Ξ	379 ± 4	379
Ξ^*	594 ± 2	614
Ω	733 ± 1	751

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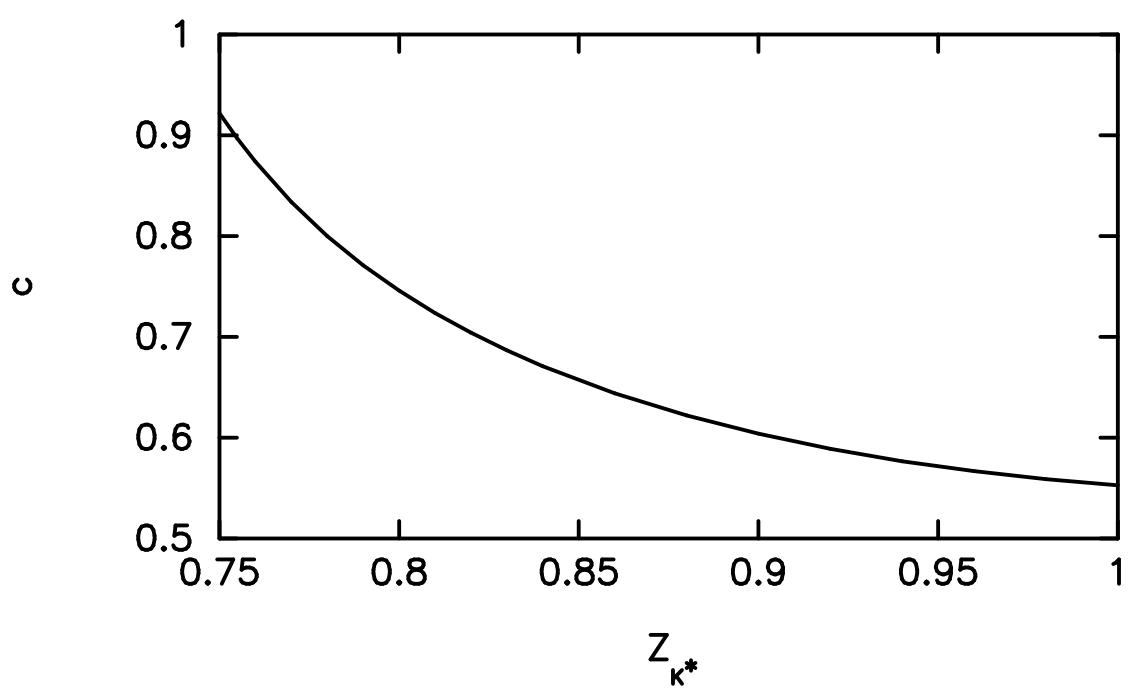
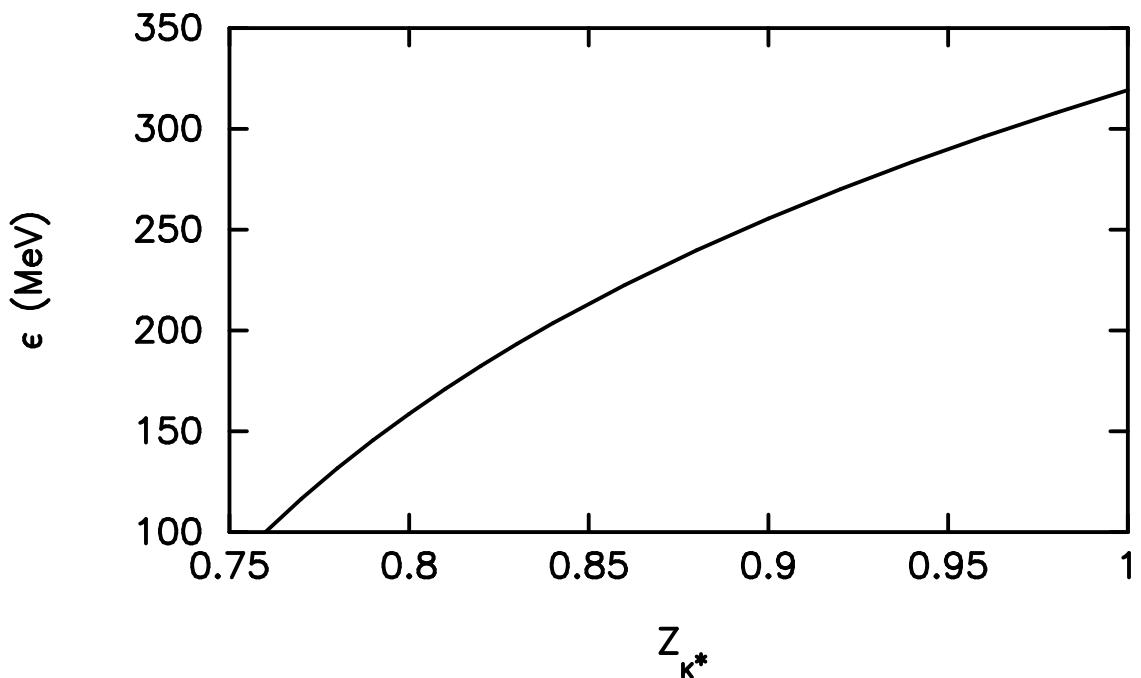


Fig.1